

Dip-hump temperature dependence of Specific Heat and Effects of Pairing Fluctuations in the Weak-coupling Side of a p -wave Interacting Fermi Gas

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We investigate the specific heat C_V at constant volume in the normal state of a p -wave interacting Fermi gas. Including p -wave pairing fluctuations within the strong-coupling theory developed by Nozières and Schmitt-Rink, we show that, in the weak-coupling side, C_V exhibits a dip-hump behavior as a function of the temperature. While the dip is associated with the pseudogap phenomenon near T_c , the hump structure is found to come from the suppression of Fermi quasiparticle scattering into a p -wave molecular state in the Fermi degenerate regime. Since the latter phenomenon does not occur in the ordinary s -wave interacting Fermi gas, it may be viewed as a characteristic phenomenon associated with a p -wave pairing interaction.

In an ultracold Fermi gas with a p -wave pairing interaction, the most characteristic phenomenon is the p -wave superfluid phase transition. Since the realization of a tunable p -wave interaction associated with a Feshbach resonance, extensive experimental^{1–11} and theoretical efforts^{12–27} have been done to realize this unconventional Fermi superfluid. However, although the formation of p -wave molecules has been reported,⁷ no one has succeeded in this exciting challenge. This is because, although a p -wave interaction is always needed to realize a p -wave superfluid Fermi gas, it also causes the dipolar relaxation,^{2, 7–9} as well as three-body loss,^{28, 29} leading to very short lifetime ($\tau_{\text{Cooper}} = 5 \sim 20$ ms) of p -wave Cooper pairs,³⁰ as well as particle loss from the system. As a result, these pairs are soon destroyed before the condensation growth ($= O(100 \text{ ms}) \gg \tau_{\text{Cooper}}$).

Because of this difficulty, in the current stage of research on a p -wave interacting Fermi gas, it is interesting to look for other characteristic phenomena that are absent in the ordinary s -wave case. Since the current experiment is only accessible to the normal state of a p -wave interacting Fermi gas, one should explore such phenomena above T_c . To avoid the serious particle loss^{2, 7–9, 28, 29} as much as possible, the weak-coupling regime away from a p -wave Feshbach resonance would be better.

In this letter, we present a phenomenon that meets the above demands, within the framework of the strong-coupling theory developed by Nozières and Schmitt-Rink. Usually, interaction effects are considered to be weak in the weak-coupling regime. However, even in this regime, we show that a p -wave pairing interaction still gives an anomalous temperature dependence of the specific heat C_V at constant volume above T_c , which is not seen in the s -wave case at all. We briefly note that the observation of this thermodynamic quantity has recently become possible in cold Fermi gas physics.³¹

We consider a one-component uniform Fermi gas with a p -wave interaction, described by the Hamiltonian,^{13, 14, 26, 27}

$$H = \sum_p \xi_p c_p^\dagger c_p - \frac{1}{2} \sum_{p, p', q} V_p(\mathbf{p}, \mathbf{p}') c_{p+q/2}^\dagger c_{-p+q/2}^\dagger c_{-p'+q/2} c_{p'+q/2}. \quad (1)$$

In this letter, we set $\hbar = k_B = 1$, and the system volume

V is taken to be unity, for simplicity. Here, c_p is an annihilation operator of a Fermi atom with the kinetic energy $\xi_p = p^2/(2m) - \mu$, measured from the Fermi chemical potential μ (where m is an atomic mass). In Eq. (1),

$$V_p(\mathbf{p}, \mathbf{p}') = - \sum_{i=x,y,z} \gamma_p^i U \gamma_{p'}^i \quad (2)$$

is an assumed p -wave pairing interaction,^{13, 14, 26, 27} where $U (> 0)$ is a coupling constant. The p -wave basis functions $\gamma_p^i = p_i F_c(\mathbf{p})$ ($i = x, y, z$) involve a cutoff function $F_c(\mathbf{p})$, which will be specified soon later. In a real ultracold Fermi gas, a p -wave interaction can be tuned by adjusting the threshold energy of a Feshbach resonance.³² However, we do not consider the detailed Feshbach mechanism in this letter, but simply treat U as a tunable parameter. As observed in a ⁴⁰K Fermi gas,³ a p -wave interaction in an ultracold Fermi gas may have a uniaxial anisotropy ($U_x \geq U_y = U_z$, where the x axis is chosen to be parallel to an external magnetic field to adjust a Feshbach resonance), because of the splitting of a p -wave Feshbach resonance by a magnetic dipole-dipole interaction. However, we ignore this, to only deal with the isotropic case ($U_x = U_y = U_z \equiv U$), for simplicity. Effects of the uniaxial anisotropy will be separately discussed in our future paper. We also ignore effects of a harmonic trap.

The cutoff function $F_c(\mathbf{p})$ in the basis function γ_p^i is to eliminate the well-known ultraviolet divergence involved in the model Hamiltonian in Eq. (1). Here, we take $F_c(\mathbf{p}) = 1/[1 + (p/p_c)^{2n}]$ with $n = 3$.²⁷ We briefly note that, while superfluid properties below T_c somehow depend on the value of n , normal state properties above T_c do not,²⁷ as far as the cutoff momentum p_0 is taken to be much larger than the Fermi momentum k_F . Although the momentum cutoff p_c and the bare p -wave interaction strength U are not observable, they are related to the observable scattering volume v , as well as the inverse effective range k_0 , as

$$\frac{4\pi v}{m} = -\frac{U}{3} \frac{1}{1 - \frac{U}{3} \sum_p \frac{p^2}{2\varepsilon_p} F_c^2(\mathbf{p})}, \quad (3)$$

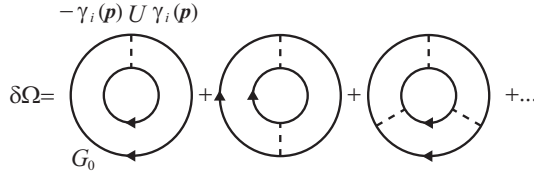


Fig. 1. p -wave fluctuation corrections $\delta\Omega$ to the thermodynamic potential Ω . The solid line and the dashed line represent, respectively, the bare Green's function $G_0(\mathbf{p}, i\omega_n) = (i\omega_n - \xi_p)^{-1}$, and a p -wave pairing interaction $V_p(\mathbf{p}, \mathbf{p}')$, where ω_n is the fermion Matsubara frequency.

$$k_0 = -\frac{4\pi}{m^2} \sum_p \frac{p^2}{2\varepsilon_p^2} F_c^2(\mathbf{p}). \quad (4)$$

We take $k_0 = -30k_F$ (where k_F is the Fermi momentum), following the experiment on a ^{40}K Fermi gas.³⁾ When we measure the strength of a p -wave interaction in terms of the inverse scattering volume v^{-1} , the weak-coupling side and the strong-coupling side are conveniently characterized as $(k_F^3 v)^{-1} \lesssim 0$ and $(k_F^3 v)^{-1} \gtrsim 0$, respectively. We briefly note that, although there is actually no clear boundary between the two regions, the Fermi chemical potential $\mu(T \simeq T_c)$ becomes negative around $(k_F^3 v)^{-1} = 0$,^{13,14)} indicating that the system gradually becomes dominated by two-body bound molecules, as one passes through $(k_F^3 v)^{-1} = 0$.

We include p -wave pairing fluctuations within the framework of the strong-coupling theory developed by Nozières and Schmitt-Rink (NSR). The advantage of this approach is that it directly evaluates fluctuation corrections $\delta\Omega$ to the thermodynamic potential $\Omega = \Omega_0 + \delta\Omega$, that are diagrammatically given in Fig. 1 (where $\Omega_0 = T \sum_p \ln[1 + e^{-\xi_p/T}]$ is the thermodynamic potential of a free Fermi gas). Once Ω is obtained, one can calculate the specific heat $C_V = (\partial E / \partial T)_{V,N}$ at constant volume, by evaluating the internal energy E from the Legendre transformation,

$$E = \Omega - T \left(\frac{\partial \Omega}{\partial T} \right)_\mu - \mu \left(\frac{\partial \Omega}{\partial \mu} \right)_T. \quad (5)$$

Summing up the diagrams in Fig. 1, we obtain,^{26,27)}

$$\delta\Omega = T \sum_{i\nu_n, \mathbf{q}} \text{Tr} \left[\ln \hat{\Gamma}(\mathbf{q}, i\nu_n) \right], \quad (6)$$

where ν_n is the boson Matsubara frequency, and

$$\hat{\Gamma}(\mathbf{q}, i\nu_n) = -\frac{U}{1 - U\hat{\Pi}(\mathbf{q}, i\nu_n)} \quad (7)$$

is the 3×3 -matrix particle-particle scattering matrix, describing fluctuations in the p -wave Cooper channel. In Eq. (7), $\hat{\Pi} = \{\Pi_{i,j}\}$ ($i, j = x, y, z$) is the lowest-order pair correlation function, given by

$$\Pi_{i,j}(\mathbf{q}, i\nu_n) = -\frac{1}{2} \sum_p \gamma_p^i \gamma_p^j \frac{1 - f(\xi_{\mathbf{p}+\frac{\mathbf{q}}{2}}) - f(\xi_{\mathbf{p}-\frac{\mathbf{q}}{2}})}{i\nu_n - \xi_{\mathbf{p}+\frac{\mathbf{q}}{2}} - \xi_{\mathbf{p}-\frac{\mathbf{q}}{2}}}. \quad (8)$$

Here, $f(x)$ is the Fermi distribution function. We briefly note that by taking z axis along the direction of \mathbf{q} , $\hat{\Pi}$, as well as $\hat{\Gamma}$ become diagonal in the present isotropic p -wave interaction above T_c .

To calculate C_V in the normal state, we need to evaluate the Fermi chemical potential $\mu(T \geq T_c)$, which is, as usual, deter-

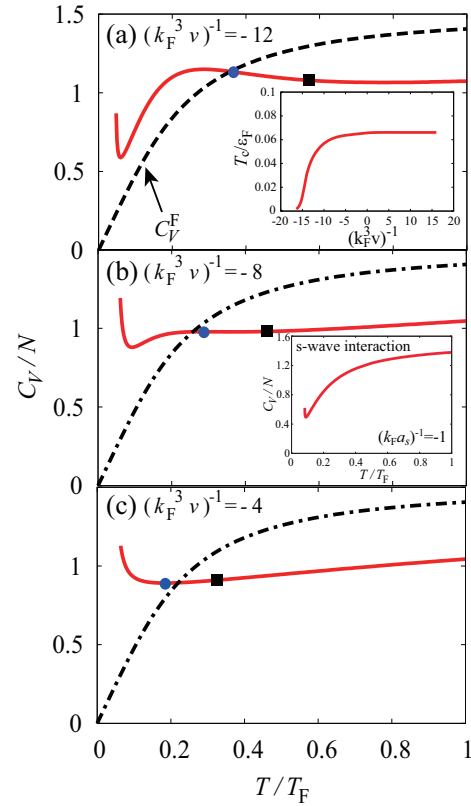


Fig. 2. (Color online) Calculated specific heat C_V at constant volume in the weak-coupling side $(k_F^3 v)^{-1} \lesssim 0$ of a one-component p -wave interacting Fermi gas above T_c . C_V^F is the specific heat in a free Fermi gas. The solid square and solid circle show the temperature at which $\mu(T) = 0$ and $\mu(T) = T$, respectively (see also Fig. 4). The inset in panel (a) shows T_c , as a function of the p -wave interaction strength. ε_F is the Fermi energy. The inset in panel (b) shows C_V in the weak-coupling regime of an s -wave interacting Fermi gas, where a_s is the s -wave scattering length.

mined from the equation for the number N of Fermi atoms,³³⁾

$$N = -\left(\frac{\partial \Omega}{\partial \mu} \right)_T = -\left(\frac{\partial \Omega_0}{\partial \mu} \right)_T - \left(\frac{\partial \delta\Omega}{\partial \mu} \right)_T. \quad (9)$$

We briefly note that, in the NSR scheme, the superfluid phase transition temperature T_c is determined from the Thouless criterion, $\Gamma^{-1}(0, 0) = 0$.^{26,27,33)}

Figure 2 shows the specific heat C_V at constant volume above T_c in the weak-coupling side of a p -wave interacting Fermi gas $((k_F^3 v)^{-1} < 0$, see also the inset in Fig. 2(a)). In panel (a), one clearly sees a dip-hump structure. Figures 2(b) and (c) show that this structure gradually becomes obscure, with increasing the interaction strength. In Fig. 2(c), while the dip still remains, the hump no longer exists. In a free Fermi gas, the specific heat C_V exhibits a monotonic temperature dependence (dashed line in Fig. 2). In addition, such a hump structure is also absent in an s -wave interacting Fermi gas, as shown in the inset in Fig. 2(b). Thus, the p -wave interaction is found to play a crucial role for the appearance of the hump structure of C_V .

To understand the behavior of C_V shown in Fig. 2, it is convenient to consider the single-particle thermal Green's function G which is consistent with the NSR theory (in the sense that it gives the same number equation as Eq. (9)), given by

$$G(\mathbf{p}, i\omega_n) = G_0(\mathbf{p}, \omega_n) + G_0(\mathbf{p}, \omega_n) \Sigma(\mathbf{p}, \omega_n) G_0(\mathbf{p}, \omega_n)$$

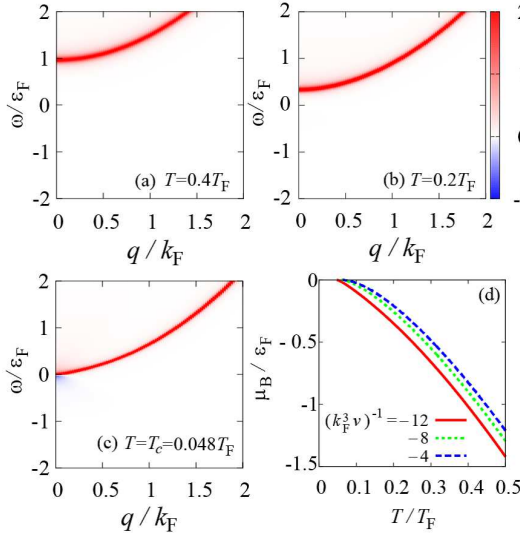


Fig. 3. (Color online) (a)-(c) Calculated spectra of the analytic continued particle-particle scattering matrix, $-\text{Im}[\Gamma(\mathbf{q}, iv_n \rightarrow \omega + i\delta)]$. The intensity is normalized by ε_F/k_F^5 . The same normalization is also used in Figs. 6(a)-(c). We set $(k_F^3 v)^{-1} = -12$. (d) Bose chemical potential μ_B , as a function of temperature.

$$\simeq \frac{1}{i\omega_n - \xi_p - \Sigma(\mathbf{p}, i\omega_n)}, \quad (10)$$

where $G_0(\mathbf{p}, i\omega_n) = [i\omega_n - \xi_p]^{-1}$ is the bare Green's function, ω_n is the fermion Matsubara frequency, and the self-energy $\Sigma(\mathbf{p}, i\omega_n)$ has the form,²⁶⁾

$$\Sigma(\mathbf{p}, i\omega) = T \sum_{\mathbf{q}, iv_n} \sum_{i=x,y,z} \gamma_{p-q/2}^i \Gamma(\mathbf{q}, iv_n) \gamma_{p-q/2}^i G_0(\mathbf{q}-\mathbf{p}, iv_n - i\omega_n). \quad (11)$$

Near T_c , since the particle-particle scattering matrix $\Gamma(\mathbf{q}, iv_n)$ is enhanced around $\mathbf{q} = \mathbf{v}_n = 0$,³³⁾ the self-energy in Eq. (11) may be approximated to²⁶⁾

$$\Sigma(\mathbf{p}, i\omega_n) \simeq -\Delta_{\text{pg}}^2(\mathbf{p}) G_0(-\mathbf{p}, -i\omega_n), \quad (12)$$

where $\Delta_{\text{pg}}^2 = -T \mathbf{p}^2 \sum_{\mathbf{q}, v_n} \Gamma(\mathbf{q}, iv_n) (\geq 0)$ is the so-called pseudogap parameter,³⁴⁾ describing a particle-hole coupling by pairing fluctuations. When one substitutes Eq. (12) into the second line in Eq. (10), the resulting Green's function has the same form as the diagonal component of the mean-field BCS Green's function in the superfluid phase as,

$$G(\mathbf{p}, i\omega_n) = -\frac{i\omega_n + \xi_p}{\omega_n^2 + \xi_p^2 + \Delta_{\text{pg}}^2(\mathbf{p})}. \quad (13)$$

Equation (13) indicates that the system near T_c has superfluid-phase-like properties, which is sometimes referred to as the pseudogap phenomenon in the literature. As a result, the entropy S is suppressed near T_c , leading to the enhancement of $C_V = T(\partial S/\partial T)_{V,N}$ as seen in Fig. 2 below the dip structure.

We note that the same pseudogap phenomenon is also seen in the s -wave case (inset in Fig. 2(b)).³⁵⁾ Thus, the enhancement of C_V near T_c and the resulting dip structure seen in Fig. 2 should be regarded as a phenomenon widely in attractively interacting Fermi gases near T_c , rather than a phenomenon peculiar to the p -wave case.

On the other hand, the key to understand the hump structure seen in Fig. 2(a) is that the relatively large magnitude of the inverse effective range $|k_c/k_-| = 30 \gg 1$ ³⁾ as

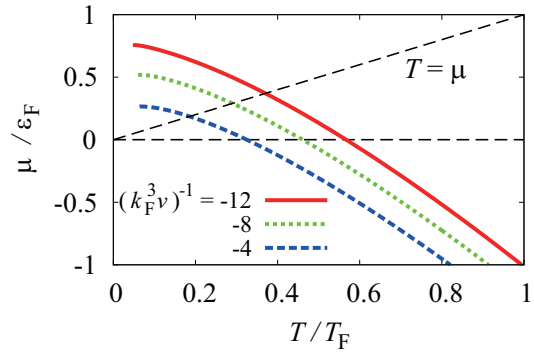


Fig. 4. (Color online) Calculated Fermi chemical potential μ , as a function of temperature. We note that, in each panel in Fig. 2, the solid square and solid circles are, respectively, the temperature at which $\mu(T) = 0$ and the temperature at which $\mu(T)$ intersects the dashed line ($\mu = T$).

well as the momentum dependence of the p -wave interaction ($V_p(\mathbf{p}, \mathbf{p}') \propto p p'$), lead to the Bose-propagator-like structure of the particle-particle scattering matrix $\Gamma(\mathbf{q}, iv_n)$ as²⁶⁾

$$\Gamma(\mathbf{q}, iv_n) \simeq \frac{Z}{iv_n - E_q}, \quad (14)$$

where $Z = 24\pi/(m^2|k_0|)$, and $E_q = \mathbf{q}^2/(4m) - \mu_B$ with $\mu_B = 2\mu - 2/(mvk_0)$. Indeed, as shown in Figs. 3(a)-(c), the spectrum $-\text{Im}[\Gamma(\mathbf{q}, iv_n \rightarrow \omega + i\delta)]$ has a sharp peak line along the molecular dispersion $\omega = E_q$ (where δ is an infinitesimally small positive number). Substituting Eq. (14) into Eq. (12), the Fermi quasiparticle damping rate, $\bar{\gamma}(\mathbf{p}, \omega) = -\text{Im}[\Sigma(\mathbf{p}, i\omega_n \rightarrow \omega + i\delta)]$, is evaluated as

$$\begin{aligned} \bar{\gamma}(\mathbf{p}, \omega) &= -\pi Z \sum_{\mathbf{q}} \left[\mathbf{p} + \frac{\mathbf{q}}{2} \right]^2 \\ &\times \left[f(\xi_{q-p}) + n_B(E_q) \right] \delta(\omega - E_q + \xi_{q-p}). \end{aligned} \quad (15)$$

Here, $n_B(x)$ is the Bose distribution function. For simplicity, when we consider $\bar{\gamma}(\mathbf{q}, \omega = 0)$ in the case of Fig. 2(a), Eq. (15) indicates that, with decreasing the temperature from the high temperature region where $\mu < 0$, the Fermi quasiparticle scattering starts to be suppressed, when μ becomes positive and most Fermi atoms are occupied below μ . This is because the constraint $\xi_{q-p} = E_q \geq |\mu_B| \sim \varepsilon_F$ (see Fig. 3(d)) coming from the δ -function in Eq. (15) becomes difficult to be satisfied. This suppression decreases the entropy S , which naturally leads to the increase of the specific heat $C_V = T(\partial S/\partial T)_{V,N}$ with decreasing the temperature. Indeed, Fig. 2(a) shows that the hump structure of C_V starts to appear around the temperature below which the chemical potential μ becomes positive (solid square). In addition, once the Fermi degeneracy is achieved, the temperature dependence of the suppression of quasiparticle scattering, as well as the temperature dependence of the suppression of the entropy S , would become weak. At such low temperatures, C_V is expected to decrease with decreasing the temperature, as in the case of a degenerate free Fermi gas (except near T_c). Indeed, the temperature at the top of the hump structure in Fig. 2(a) is close to the temperature which satisfies $T = \mu(T)$ (solid circle).

With increasing the interaction strength, the pseudogap regime where C_V is enhanced becomes wide. At the same time, the temperature below which the above-mentioned

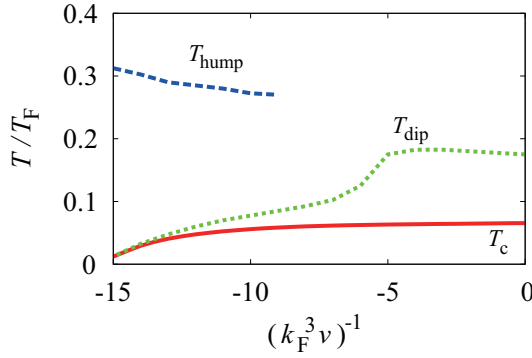


Fig. 5. Two characteristic temperatures T_{hump} and T_{dip} that are determined as the temperature at the top of the hump structure and the temperature at the dip of the specific heat C_V , respectively. T_c is the p -wave superfluid phase transition temperature.

mechanism of the hump structure works becomes low, because the temperature at which μ changes its sign becomes low, as shown in Fig. 4. These naturally explain why the hump structure of C_V gradually disappears, as one approaches the intermediate coupling regime from the weak-coupling side. When we plot the temperature T_{hump} at the top of the hump structure, we numerically confirm that T_{hump} disappears at $(k_F^3 v)^{-1} \simeq -8$, as shown in Fig. 5. On the other hand, when we plot the temperature T_{dip} at the dip, we find in Fig. 5 that it continues to exist even in the intermediate coupling regime.

We note that all the p -wave pairing interactions do not give the dip-hump behavior of C_V . Indeed, when we only retain the p -wave symmetry but ignore the $|\mathbf{p}|$ -dependence of the prefactor of $F_c(\mathbf{p})$ in the basis function as $\gamma_p^i = (\mathbf{p}/|\mathbf{p}|)F_c(\mathbf{p})$, C_V does not exhibit the hump structure, as shown in Fig. 6(d). In the case of $\gamma_p^i = \mathbf{p}F_c(\mathbf{p})$, the p -wave interaction is enhanced in the high-energy region by the factor $\mathbf{p} \cdot \mathbf{p}'$ in $V_p(\mathbf{p}, \mathbf{p}')$ in Eq. (2), which also enhances the molecular character of pairing fluctuations, leading to the sharp peak line along E_q in Figs. 3(a)-(c). In contrast, this effect is absent in the case of $\gamma_p^i = (\mathbf{p}/|\mathbf{p}|)F_c(\mathbf{p})$. Thus, the spectrum of the particle-particle scattering matrix $\Gamma(\mathbf{q}, i\nu_n \rightarrow \omega + i\delta)$ spreads out, as shown in Figs. 6(a)-(c). In this case, the constraint coming from the δ -function in Eq. (15) is not obtained. As a result, even when the Fermi chemical potential μ becomes positive and the Fermi degeneracy starts to occur, the suppression of the entropy S is not remarkable enough to cause the enhancement of C_V .

Since the ordinary contact-type s -wave pairing interaction also does not have a factor which enhances the interaction strength in the high-energy region, the situation is similar to the case of $\gamma_p^i = (\mathbf{p}/|\mathbf{p}|)F_c(\mathbf{p})$. Thus, the spectrum of $\Gamma(\mathbf{q}, i\nu_n \rightarrow \omega + i\delta)$ also spreads out in this case, although we do not explicitly show the result here. This explains the reason for the absence of the hump structure in the s -wave case shown in the inset in Fig. 2(b).

Strictly speaking, the above-mentioned mechanism for the hump structure would also work for other unconventional pairing interactions, such as a d -wave one. Thus, the observation of a dip-hump structure of C_V does not immediately mean that a p -wave interaction play an essential role. However, the advantage of an ultracold Fermi gas is that one can unambiguously pick up the symmetry of pairing interaction by using a Feshbach resonance. Thus, when a hump structure

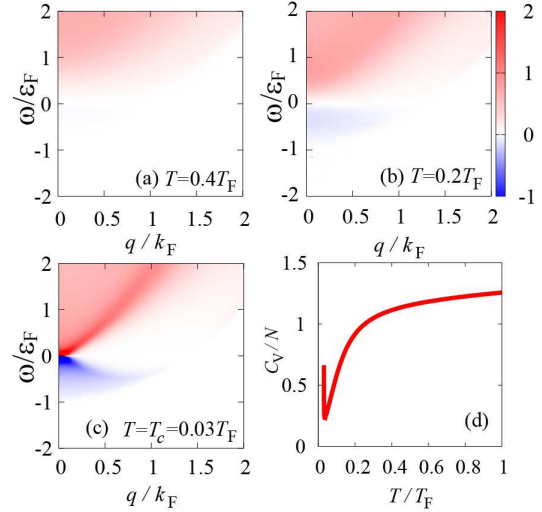


Fig. 6. (Color online) (a)-(c) Calculated spectrum of the analytic continued particle-particle scattering matrix, $-\text{Im}[\Gamma(\mathbf{q}, i\nu_n \rightarrow \omega + i\delta)]$, when the basis function $\gamma_p^i = (\mathbf{p}/|\mathbf{p}|)F_c(\mathbf{p})$ is used. We take $(k_F^3 v)^{-1} = -1$. The panel (d) shows C_V .

of C_V is observed near a p -wave Feshbach resonance, we may actually conclude that it comes from the $|\mathbf{p}|$ -dependence of a p -wave pairing interaction associated with the p -wave Feshbach resonance.

To summarize, we have theoretically discussed the specific heat C_V at constant volume in the weak-coupling side of a p -wave interacting Fermi gas. Within the framework of the strong-coupling theory developed by Nozières and Schmitt-Rink, we showed that a p -wave interaction causes a dip-hump behavior of C_V as a function of the temperature above T_c . While the dip structure, which is related to the pseudogap phenomenon, also appears in an s -wave interacting Fermi gas, the hump structure does not. Thus, the latter may be viewed as a characteristic phenomenon in a p -wave interacting Fermi gas. Since the achievement of the p -wave superfluid phase, which is the most characteristic phenomenon in a p -wave interacting Fermi gas, seems still difficult in the current stage of cold Fermi gas physics, the observation of this normal-state phenomenon expected in the weak-coupling side of a p -wave interacting Fermi gas would be helpful to clarify how this system is qualitatively different from the ordinary s -wave system that has extensively been studied in this field.

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